Constructive use of holographic projections

Dedicated to Klaus Fredenhagen on the occasion of his 60th birthday

Bert Schroer
CBPF, Rua Dr. Xavier Sigaud 150
22290-180 Rio de Janeiro, Brazil
and Institut fuer Theoretische Physik der FU Berlin, Germany

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Abstract

Revisiting the old problem of existence of interacting models of QFT with new conceptual ideas and mathematical tools, one arrives at a novel view about the nature of QFT. The recent success of algebraic methods in establishing the existence of factorizing models suggests new directions for a more intrinsic constructive approach beyond Lagrangian quantization. Holographic projection simplifies certain properties of the bulk theory and hence is a promising new tool for these new attempts.

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1 Historical background and present motivations for holography

No other theory in the history of physics has been able to cover such a wide range of phenomena with impressive precision as QFT. However its amazing predictive power stands in a worrisome contrast to its weak ontological status. In fact QFT is the only theory of immense epistemic strength which, even after more than 80 years, remained on shaky mathematical and conceptual grounds. Unlike any other area of physics, including QM, there are simply no interesting mathematically controllable interacting models which would show that the underlying principles remain free of internal contradictions in the presence of interactions. The faith in e.g. the Standard Model is based primarily on its perturbative descriptive power; outside the perturbative domain there are more doubts than supporting arguments.

The suspicion that this state of affairs may be related to the conceptual and mathematical weakness of the method of Lagrangian quantization rather then a shortcoming indicating an inconsistency of the underlying principles in the presence of interactions can be traced back to its discoverer Pascual Jordan. It

certainly was behind all later attempts of e.g. Arthur Wightman and Rudolf Haag to find a more autonomous setting away from the quantization parallelism with classical theories which culminated in Wightman's axiomatic setting in terms of vacuum correlation functions and the Haag-Kastler theory of nets of operator algebras.

The distance of such conceptual improvements to the applied world of calculations has unfortunately persisted. Nowhere is the contrast between computational triumph and conceptual misery more visible than in renormalized perturbation theory which has remained our only means to explore the standard model. Most particle physicists have a working knowledge of perturbation theory and at least some of them took notice of the fact that, although the renormalized perturbative series can be shown to diverge and that in certain cases these divergent series are Borel resummable. Here I will add some more comments without going into details.

The Borel re-summability property unfortunately does not lead to an existence proof; the correct mathematical statement in this situation is that if the existence can be established[1] by nonperturbative method then the Borel-resummed series would indeed acquire an asymptotic convergence status with respect to the solution, and one would for the first time be allowed to celebrate the numerical success as having a solid ontological basis [2]. But the whole issue of model existence attained the status of an unpleasant fact, something which is often kept away from newcomers, so that as a result there is a certain danger to confuse the existence of a model with the ability to write down a Lagrangian or a functional integral and apply some computational recipe.

Fortunately important but unfashionable problems in particle physics never disappear completely. Even if they have been left on the wayside as "unstringy", "unsupersymmetrizable" or too far removed from the "Holy Grail of a TOE" and therefore not really career-improving, there will be always be individuals who return to them with new ideas.

Indeed there has been some recent progress about the aforementioned existence problem from a quite unexpected direction. Within the setting of d=1+1 factorizing models the use of modular operator theory has led to a control over phase space degrees of freedom which in turn paved the way to an existence proof. Those models are distinguished by their simple generators for the wedge-localized algebra [4]; in fact these generators turned out to possess Fourier-transforms with mass-shell creation/annihilation operators which are only slightly more complicated than free fields. An important additional idea on the way to an existence proof is the issue of the cardinality of degrees of freedom. In the form of the phase space in QFT as opposed to QM this issue goes back to the 60s [1] and underwent several refinements [2] (a sketch of the history can be found in [3]).

[1]The existence for models with a finite wave-function renormalization constant has been established in the early 60s and this situation has not changed up to recently. The old results only include superrenormalizable models whereas the new criterion is not related to short-distance restrictions but rather requires a certain phase space behavior (modular nuclearity).

[2]This is actually the present situation for the class of d=1+1 factorizing models [5].

The remaining problem was to show that the simplicity of the wedge generators led to a "tame" phase space behavior which guaranties the nontriviality as well as the additional expected properties of the double cone localized algebras obtained as intersections of wedge-localized algebras [5]. Although these models have no particle creation through on-shell scattering, they exhibit the full infinite vacuum polarization clouds upon sharpening the localization from wedges to compact spacetime regions as e.g. double cones [6]. Their simplicity is only manifest in the existence of simple wedge generators; for compact localization regions their complicated infinite vacuum polarization clouds are not simpler than in other QFT.

Similar simple-minded Ansätze for wedge algebras in higher dimensions cannot work since interactions which lead to nontrivial elastic scattering without also causing particle creation cannot exist; such a No-Go theorem for 4-dim. QFT was established already in [7]. Nevertheless it is quite interesting to note that even if with such a simple-minded Ansatz for wedge generators in higher dimensions one does not get to compactly localized local observables, one can in some cases go to certain subwedge intersections [8][9] before the increase in localization leads to trivial algebras.

Whereas in the Lagrangian approach one starts with local fields and their correlations and moves afterwards to less local objects as global charges, incoming fields[3] etc., the modular localization approach goes the opposite way i.e. one starts from the wedge region (the best compromise between particles and fields) which is most close to the particle mass-shell the S-matrix and then works one's way down. The pointlike local fields only appear at the very end and play the role of *coordinatizing generators* of the double cone algebras for arbitrary small sizes.

Nonlocal models are automatically "noncommutative" in the sense that the the maximal commutativity of massive theories allowed by the principles of QFT, namely spacelike commutativity, is weakened by allowing various degrees of violations of spacelike commutativity. In this context the non-commutativity associated with the deformation of the product to a star-product using the Weyl-Moyal formalism is only a very special (but very popular) case. The motivation for studying non-commutative QFT for its own sake comes from string theory, and one should not expect this motivation to be better than for string theory itself.

My motivation for having being interested in noncommutative theory during the last decade comes from the observation that non-commutative fields can have *simpler properties than commutative ones*. More concretely: complicated two-dimensional local theories may lead to wedge-localized algebras which are generated by non-commutative fields where the latter only fulfill the much weaker wedge-locality (see above). Whereas in d=1+1 such constructions [4] may lead via algebraic intersections to nontrivial, nonperturbative local fields, it is known that in higher dimensions this simple kind of wedge generating field

[3]Incoming/outgoing free fields are only local with respect to themselves. The physically relevant notion of locality is *relative locality to the interacting fields.* If incoming fields are relatively local/almost local, the theory has no interactions.

without vacuum polarization is not available. But interestingly enough on can improve the wedge localization somewhat [10] before the further sharpening of localization via algebraic intersections ends in trivial algebras.

These recent developments combine the useful part of the history of S-matrix theory and formfactors with very new conceptual inroads into QFT (modular localization, phase space properties of LQP). The idea to divide the difficult full problem into a collection of simpler smaller ones is also at the root of the various forms of the holography of the two subsequent sections.

The predecessor of lightfront holography was the so-called "lightcone quantization" which started in the early 70s; it was designed to focus on short-distances and forget temporarily about the rest. The idea to work with fields which are associated to the lightfront $x_- = 0$ (not the light cone which is $x^2 = 0$) as a submanifold in Minkowski spacetime looked very promising but unfortunately the connection with the original problem of analyzing the local theory in the bulk was never addressed and as the misleading name "lightcone quantization" reveals, the approach was considered as a different quantization rather then a different method for looking at the same local QFT in Minkowski spacetime. It is not really necessary to continue a seperate criticism of "lightcone quantization" because its shortcomings will be become obvious after the presentation of lightfront holography (more generally *holography onto null-surfaces*).

Whereas the more elaborate and potentially more important lightfront holography has not led to heated discussions, the controversial potential of the simpler AdS-CFT holography had been enormous and to the degree that it contains interesting messages which increase our scientific understanding it will be presented in these notes.

Since all subjects have been treated in the existing literature, our presentation should be viewed ass a guide through the literature with occasionally additional and (hopefully) helpful remarks.

2 Lightfront holography, holography on null-surfaces and the origin of the area law

Free fields offer a nice introduction into the bulk-holography relation which, despite its simplicity, remains conceptually non-trivial.

We seek generating fields A_{LF} for the lightfront algebra $\mathcal{A}(LF)$ by following the formal prescription $x_- = 0$ of the old "lightfront approach" [11]. Using the abbreviation $x_\pm = x^0 \pm x^3,\ p_\pm = p^0 + p^3 \simeq e^{\mp\theta},$ with θ the momentum space rapidity :

$$A_{LF}(x_+, x_\perp) := A(x)|_{x_-=0} \simeq \int \left(e^{i(p_-(\theta)x_+ + ip_\perp x_\perp} a^*(\theta, p_\perp) d\theta dp_\perp + h.c. \right) \quad (1)$$

$$\left\langle \partial_{x_+} A_{LF}(x_+, x_\perp) \partial_{x\prime_+} A_{LF}(x'_+, x'_\perp) \right\rangle \simeq \frac{1}{\left(x_+ - x'_+ + i\varepsilon\right)^2} \cdot \delta(x_\perp - x'_\perp)$$

$$\left[\partial_{x_+} A_{LF}(x_+, x_\perp), \partial_{x\prime_+} A_{LF}(x'_+, x'_\perp) \right] \simeq \delta'(x_+ - x'_+) \delta(x_\perp - x'_\perp)$$

The justification for this formal manipulation[4] follow from the fact that the equivalence class of test function $[f]$, which have the same mass shell restriction $\tilde{f}|_{H_m}$ to the mass hyperboloid of mass m, is mapped to a unique test function f_{LF} which "lives"on the lightfront [12][13]. It only takes the margin of a newspaper to verify the identity $A(f) = A([f]) = A_{LF}(f_{LF})$. This identity does not mean that the A_{LF} generator can be used to describe the local substructure in the bulk. The inversion involves an equivalence class and does not distinguish an individual test-function in the bulk; in fact a finitely localized test function $f(x_+, x_\perp)$ on LF corresponds to a de-localized subspace in the bulk. Using an intuitive metaphoric language one may say that a strict localization on LF corresponds to a fuzzy localization in the bulk and vice versa. Hence the pointwise use of the LF generators enforces the LF localization and the only wedge-localized operators which can be directly obtained as smeared A_{LF} fields have a noncompact extension within a wedge whose causal horizon is on LF. Nevertheless there is equality between the two operator algebras associated to the bulk W and its (upper) horizon ∂W

$$\mathcal{A}(W) = \mathcal{A}(H(W)) \subset \mathcal{A}(LF) = B(H) \quad (2)$$

These operator algebras are the von Neumann closures of the Weyl algebras generated by the smeared fields A and A_{LF} and it is only in the sense of this closure (or by forming the double commutant) that the equality holds. Quantum field theorists are used to deal with single operators. Therefore the knowledge about the equality of algebras without being able to say which operators are localized in subregion is somewhat unaccustomed. As will be explained later on, the finer localization properties in the algebraic setting can be recovered by taking suitable intersections of wedge algebras i.e. the structure of the family of all wedge algebras determines whether the local algebras are nontrivial and in case they are permits to compute the local net which contains all informations about the particular model.

This idea of taking the holographic projection of individual bulk fields can be generalized to composites of free fields (as e.g. the stress-energy tensor). In order to avoid lengthy discussions about how to interpret logarithmic chiral two-point functions in terms of restricted test functions[5] we work restrict our

[4] We took the derivatives for technical reasons (in order to write the formulas without test functions).

[5] This is a well-understood problem of chiral fields of zero scale dimension which is not directly related to holography.

attention to Wick-composites of $\partial_{x_+} A_{LF}(x_+, x_\perp)$

$$\left[B_{LF}(x_+, x_\perp), C_{LF}(x'_+, x'_\perp)\right] = \sum_{l=0}^{m} \delta^l(x_\perp - x'_\perp) \sum_{k(l)=0}^{n(l)} \delta^{k(l)}(x_+ - x'_+) D_{LF}^{(k(l))}(x_+, x_\perp) \tag{3}$$

where the dimensions of the composites $D_{LF}^{(k(l))}$ together with the degrees of the derivatives of the delta functions obey the standard rule of scale dimensional conservation. In the commutator the transverse and the longitudinal part both appear with delta functions and their derivatives yet there is a very important structural difference which shows up in the correlation functions. To understand this point we look at the second line in (1). The longitudinal (=lightlike) delta-functions carries the chiral vacuum polarization the transverse part consists only of products of delta functions as if it would come from a product of correlation functions of nonrelativistic Schroedinger creation/annihilation operators $\psi^*(x_\perp)$, $\psi(x_\perp)$. In other words the LF-fields which feature in this extended chiral theory are *chimera between QFT and QM*; they have one leg in QFT and n-2 legs in QM with the "chimeric vacuum" being partially a (transverse) factorizing quantum mechanical state of "nothingness" (the Buddhist nirvana) and partially the longitudinally particle-antiparticle polarized LQP vacuum state of "virtually everything" (the Abrahamic heaven).

Upon lightlike localization of LF to (in the present case) ∂W (or to a longitudinal interval) the vacuum on $\mathcal{A}(\partial W)$ becomes a radiating KMS thermal state with nonvanishing localization-entropy [13][14]. In case of interacting fields there is no change with respect to the absence of transverse vacuum polarization, but unlike the free case the global algebra $\mathcal{A}(LF)$ or the semi-global algebra $\mathcal{A}(\partial W)$ is generally bigger than the algebra one obtains from the globalization using compactly localized subalgebras, i.e. $\overline{\cup_{O\subset LF}\mathcal{A}_{LF}(\mathcal{O})} \subset \mathcal{A}(LF)$, $\mathcal{O} \subset LF$. We will return to this point at a more opportune moment.

The aforementioned "chimeric" behavior of the vacuum is related in a profound way to the conceptual distinctions between QM and QFT [16]. Whereas transversely the vacuum is tensor-factorizing with respect to the Born localization and therefore leads to the standard quantum mechanical concepts of entanglement and the related information theoretical (cold) entropy, the entanglement from restricting the vacuum to an algebra associated with an interval in lightray direction is a thermal KMS state with a genuine thermodynamic entropy. Instead of the standard quantum mechanical dichotomy between pure and entangled restricted states there are simply no pure states at all. All states on sharply localized operator algebras are highly mixed and the restriction of global particle states (including the vacuum) to the W-horizon $\mathcal{A}(\partial W)$ results in KMS thermal states. This is the result of the different nature of localized algebras in QFT from localized algebras in QM [16].

Therefore if one wants to use the terminology "entanglement" in QFT one should be aware that one is dealing with a totally intrinsic very strong form of entanglement: *all physically distinguished global pure states* (in particular finite energy states in particular the vacuum) *upon restriction to a localized algebra*

become intrinsically entangled and unlike in QM there is no local operation which disentangles.

Whereas the cold (information theoretic) entanglement is often linked to the uncertainty relation of QM, the raison d'etre behind the "hot" entanglement is the phenomenon of vacuum polarization resulting from localization in quantum theories with a maximal velocity. The transverse tensor factorization restricts the Reeh-Schlieder theorem (also known as the "state-operator relation"). For a longitudinal strip (st) on LF of a finite transverse extension the LF algebra tensor-factorizes together with the Hilbert space $H = H_{st} \otimes H_{st\perp}$ and the H_{st} projected form of the Reeh-Schlieder theorem for a subalgebra localized within the strip continues to be valid.

This concept of *transverse extended chiral fields* can also be axiomatically formulated for interacting fields independently of whether those objects result from a bulk theory via holographic projection or whether one wants to study QFT on (non-hyperbolic) null-surfaces. These "lightfront fields" share some important properties with chiral fields. In both cases subalgebras localized on subregions lead to a *geometric modular theory,* whereas in the bulk this property is restricted to wedge algebras. Furthermore in both cases the symmetry groups are infinite dimensional; in chiral theories the largest possible group is (after compactification) Diff($\dot{R}$), whereas the transverse extended version admits besides these pure lighlike symmetries also $x_\perp$-x_+ mixing ($x_\perp$-dependent) symmetry transformations which leave the commutation structure invariant.

There is one note of caution, unlike those conformal QFTs which arise as chiral projections from 2-dimensional conformal QFT, the extended chiral models of QFT on the lightfront which result from holography do not come with a stress-energy tensor and hence the diffeomorphism invariance beyond the Moebius invariance (for which one gets from modular invariance, no energy momentum tensor needed) is not automatic. This leads to the interesting question if there are concepts which permit to incorporate also the diffeomorphisms beyond the Moebius transformations into a modular setting, a problem which will not be pursuit here.

We have formulated the algebraic structure of holographic projected fields for bosonic fields, but it should be obvious to the reader that a generalization to Fermi fields is straightforward. Lightfront holography is consistent with the fact that except for d=1+1 there are no operators which "live" on a lightray since the presence of the quantum mechanical transverse delta function prevents such a possibility i.e. only after transverse averaging with test functions does one get to (unbounded) operators.

It is an interesting question whether a direct "holographic projection" of *interacting* pointlike bulk fields into lightfront fields analog to (1) can be formulated, thus avoiding the algebraic steps starting with wedge algebra. The important formula which led to the lightfront generators is the *mass shell representation* of the free field; if we would have performed the $x_- = 0$ limit in the two point function the result would diverge. This suggests that we should start from the so-called Glaser-Lehmann-Zimmermann (GLZ) representation [17] which is an on-shell representation in terms an infinite series of integrals

involving the incoming particle creation/annihilation operators

$$A(x) = \sum \frac{1}{n!} \int dx_1 ... \int dx_n \ a(x; x_1, ...x_n) : A_{in}(x_1) ....A(x_n) : \qquad (4)$$
$$A(x) = \sum \frac{1}{n!} \int_{H_m} dp_1 ... \int_{H_m} dp_n \ e^{ix(\sum p_i)} \tilde{a}(p_1, ...p_n) : \tilde{A}(p_1) ....\tilde{A}(p_n) :$$
$$A(x)_{LF} = A(x)_{x_- = 0}$$

in which the coefficient functions $a(x; x_1, ...x_n)$ are *retarded functions*. The second line shows that only the mass-shell restriction of these functions matter; the momentum space integration goes over the entire mass-shell and the two components of the mass hyperboloid H_m are associated with the annihilation/creation part of the Fourier transform of the incoming field. These mass-shell restrictions of the retarded coefficient functions are related to multi-particle formfactors of the field A. Clearly we can take $x_- = 0$ in this on-shell representation without apparently creating any problems in addition to the possibly bad convergence properties of such series (with or without the lightfront restriction) which they had from the start. The use of the on-shell representation (4) is essential, doing this directly in the Wightman functions would lead to meaningless divergences, as we already noticed in the free field case.

Such GLZ formulas amount to a representation of a local field in terms of other local fields in which the *relation between the two sets* of fields is very *nonlocal*. Hence this procedure is less intuitive than the algebraic method based on relative commutants and intersections of algebras. The use of a GLZ series also goes in some sense against the spirit of holography which is to simplify certain aspects[6] in order to facilitate the solution of certain properties of the theory (i.e. to preserve the original aim of the ill-defined lightcone quantization), whereas to arrive at GLZ representations one must already have solved the on-shell aspects of the model (i.e. know all its formfactors) before applying holography.

Nevertheless, in those cases where one has explicit knowledge of formfactors, as in the case of 2-dim. factorizing models mentioned in the previous section, this knowledge can be used to calculate the scaling dimensions of their associated holographic fields A_{LF}. These fields lead to more general plektonic (braid group) commutation relations which replace the bosonic relations of transverse extended chiral observables (3). We refer to [15] in which the holographic scaling dimensions for several fields in factorizing models will be calculated, including the Ising model for which an exact determination of the scaling dimension of the order field is possible. Although the holographic dimensions agree with those from the short distance analysis (which have been previously calculated in [18]), the conceptual status of holography is quite different from that of critical universality classes. The former is an exact relation between a 2-dim. factorizing model (change of the spacetime ordering of a given bulk theory) whereas the

[6]Those aspects for which holography does not simplify include particle and scattering aspects.

latter is a passing to a different QFT in the same universality class. The mentioned exact result in the case of the Ising model strengthens the hope that GLZ representations and the closely related expansions of local fields in terms of wedge algebra generating on-shell operators [15] have a better convergence status than perturbative series.

By far the conceptually and mathematically cleanest way to pass from the bulk to the lightfront is in terms of nets of operator algebras via modular theory. This method requires to start from algebras in "standard position" i.e. a pair $(\mathcal{A},\Omega)$ such that the operator algebra $\mathcal{A}$ acts cyclically on the state vector Ω i.e. $\overline{\mathcal{A}\Omega}=H$ and has no annihilators i.e. $A\Omega=0 \curvearrowright A=0$. According to the Reeh-Schlieder theorem any localized algebra $\mathcal{A}(\mathcal{O})$ forms a standard pair $(\mathcal{A}(\mathcal{O}),\Omega)$ with respect to the vacuum Ω and the best starting point for the lightfront holography is a wedge algebra since the (upper) causal horizon ∂W of the wedge W is already half the lightfront. The crux of the matter is the construction of the local substructure on ∂W. The local resolution in longitudinal (lightray) direction is done as follows.

Let W be the x_0-x_3 wedge in Minkowski spacetime which is left invariant by the x_0-x_3 Lorentz-boosts. Consider a family of wedges W_a which are obtained by sliding the W along the $x_+=x_0+x_3$ lightray by a lightlike translation $a>0$ into itself. The set of spacetime points on LF consisting of those points on ∂W_a which are spacelike to the interior of W_b for $b>a$ is denoted by $\partial W_{a,b}$; it consists of points $x_+\in(a,b)$ with an unlimited transverse part $x_\perp\in R^2$. These regions are two-sided transverse slabs on LF.

To get to intersections of finite size one may "tilt" these slabs by the action of certain subgroups in $\mathcal{G}$ which change the transverse directions. Using the 2-parametric subgroup $\mathcal{G}_2$ of $\mathcal{G}$ which is the restriction to LF of the two "translations" in the Wigner little group (i.e. the subgroup fixing the lightray in LF), it is easy to see that this is achieved by forming intersections with G_2-transformed slabs $\partial W_{a,b}$

$$\partial W_{a,b}\cap g(\partial W_{a,b}),\ \ g\in\mathcal{G}_2 \tag{5}$$

By continuing with forming intersections and unions, one can get to finite convex regions $\mathcal{O}$ of a quite general shape.

The local net on the lightfront is the collection of all local algebras $\mathcal{A}(\mathcal{O})$, $\mathcal{O}\subset LF$ and as usual their weak closure is the global algebra $\mathcal{A}_{LF}$. For interacting systems the global lightfront algebra is generally expected to be smaller than the bulk, in particular one expects

$$\begin{aligned}&\mathcal{A}_{LF}(\partial W)\subset\mathcal{A}(\partial W)=\mathcal{A}(W) \\ &\mathcal{A}_{LF}(\partial W)=\cup_{\mathcal{O}\subset\partial W}\mathcal{A}_{LF}(\mathcal{O}),\ \mathcal{A}(W)=\cup_{\mathcal{C}\subset W}\mathcal{A}(\mathcal{C})\end{aligned} \tag{6}$$

where the semi-global algebras are formed with the localization concept of their relative nets as indicated in the second line. The smaller left hand side accounts for the fact that the formation of relative commutants as $\mathcal{A}(\partial W_{a,b})$ may not maintain the standardness of the algebra because $\overline{\cup_{a,b}\mathcal{A}(\partial W_{a,b})\Omega}\subsetneqq H$. In

that case the globalization of the algebraic holography only captures a global (i.e. not localized) subalgebra of the global bulk and one could ask whether the pointlike procedure using the GLZ representation leads to generating fields which generate a bigger algebra. gives more. The answer is positive since also (bosonic) fields with anomalous short distance dimensions will pass the projective holography and become anyonic field on the lightray[7] On the other hand algebraic holography filters out bosonic fields which define the chiral obervables. These chiral observables have a DHR superselection theory. This leads to the obvious conjecture

$$Alg\{proj\ hol\} \subseteq Alg\{DHR\} \tag{7}$$

Here the left hand side denotes the algebra generated by applying projective holography to the pointlike bulk fields and the reight hand side is the smallest algebra which contains all DHR superselection sectors of the LF observable (extended chiral) algebra which resulted from algebraic holography.

It is worthwhile to emphasize that the connection between the operator algebraic and the pointlike prescription is much easier on LF than in the bulk. In the presence of conformal symmetries one has the results of Joerss [19]; looking at his theorems in the chiral setting an adaptation to the transverse extended chiral theories on LF should be straightforward. For consistency reasons such fields must fulfill (3) I hope to come back to this issue in a different context.

One motivation for being interested in lightfront holography is that it is expected to helpful in dividing the complicated problem of classifying and constructing QFTs according to intrinsic principles into several less complicated steps. In the case of d=1+1 factorizing models one does not need this holographic projection onto a chiral theory on the lightray for the mere existence proof. But e.g. for the determination of the spectrum of the short distance scale dimension, it is only holography and not the critical limit which permits to maintain the original Hilbert space setting. It is precisely this property which makes it potentially interesting for structural investigations and actual constructions of higher dimensional QFT.

Now we are well-prepared to address the main point of this section: the area law for localization entropy which follows from the absence of transverse vacuum polarization. Since this point does not depend on most of the above technicalities, it may be helpful to the reader to present the conceptual mathematical origin of this unique[8] tensor-factorization property. The relevant theorem goes back to Borchers [20] and can be stated as follows. Let $\mathcal{A}_i \subset B(H)$, $i = 1, 2$ be two operator algebras with $[\mathcal{A}_1, U(a)\mathcal{A}_2 U(a)^*] = 0\ \forall a$ and $U(a)$ a translation with *nonnegative* generator which fulfills the cluster factorization property (i.e. asymptotic factorization in correlation functions for infinitely large cluster separations) with respect to a unique $U(a)$-invariant state vector Ω[9]. It then follows

[7] The standard Boson-Fermion statistics refers to spacelike distances and the lightlike statistics resulting from projective holography is determined by the anomalous short distance dimensions of the bulk field and not by their statistics.

[8] Holography on null-surfaces is the only context in which a quantum mechanical structure enters a field theoretic setting.

[9] Locality in both directions shows that the lightlike translates $\langle \Omega \,|AU(a)B|\, \Omega\rangle$ are bound-

that the two algebras tensor factorize in the sense $\mathcal{A}_1 \vee \mathcal{A}_2 = \mathcal{A}_1 \otimes \mathcal{A}_2$ where the left hand side denotes the joint operator algebra.

In the case at hand the tensor factorization follows as soon as the open regions $\mathcal{O}_i \subset LF$ in $\mathcal{A}(\mathcal{O}_i)$ $i = 1, 2$ have no transverse overlap. The lightlike cluster factorization is weaker (only a power law) than its better known spacelike counterpart, but as a result of the analytic properties following from the *non-negative generator of lightlike translations* it enforces the asymptotic factorization to be valid at all distances. The resulting transverse factorization implies the transverse additivity of extensive quantities as energy and entropy and their behavior in lightray direction can then be calculated in terns of the associated auxiliary chiral theory. a well-known property for spacelike separations.

This result [13][14] of the transverse factorization may be summarized as follows

1. The system of LF subalgebras $\{\mathcal{A}(\mathcal{O})\}_{\mathcal{O} \subset LF}$ tensor-factorizes transversely with the vacuum being free of transverse entanglement

$$\mathcal{A}(\mathcal{O}_1 \cup \mathcal{O}_2) = \mathcal{A}(\mathcal{O}_1) \otimes \mathcal{A}(\mathcal{O}_2),\ (\mathcal{O}_1)_\perp \cap (\mathcal{O}_2)_\perp = \emptyset \qquad (8)$$
$$\langle \Omega \left| \mathcal{A}(\mathcal{O}_1) \otimes \mathcal{A}(\mathcal{O}_2) \right| \Omega \rangle = \langle \Omega \left| \mathcal{A}(\mathcal{O}_1) \right| \Omega \rangle \langle \Omega \right| \mathcal{A}(\mathcal{O}_2) \left| \Omega \right\rangle$$

2. Extensive properties as entropy and energy on LF are proportional to the extension of the transverse area.

3. The area density of localization-entropy in the vacuum state for a system with sharp localization on LF diverges logarithmically

$$s_{loc} = \lim_{\varepsilon \to 0} \frac{c}{6} \left| ln\varepsilon \right| + ... \qquad (9)$$

   where ε is the size of the interval of "fuzziness" of the boundary in the lightray direction which one has to allow in order for the vacuum polarization cloud to attenuate and the proportionality constant c is (at least in typical examples) the central extension parameter of the Witt-Virasoro algebra.

The following comments about these results are helpful in order to appreciate some of the physical consequences as well as extensions to more general null-surfaces.

As the volume divergence of the energy/entropy in a heat bath thermal system results from the thermodynamic limit of a sequence of boxed systems in a Gibbs states, the logarithmic divergence in the vacuum polarization attenuation distance ε plays an analogous role in the approximation of the semiinfinitely extended ∂W by sequences of algebras whose localization regions approach ∂W from the inside. In both cases the limiting algebras are monads whereas the approximands are type I analogs of the "box quantization" algebras. In fact in the

ary values of entire functions and the cluster property together with Liouville's theorem gives the factorization.

present conformal context the relation between the standard heat bath thermodynamic limit and the limit of vanishing attenuation length for the localization-caused vacuum polarization cloud really gord beyond an analogy and becomes an isomorphism.

This surprising result is based on two facts [13][14]. On the one hand conformal theories come with a natural covariant "box" approximation of the thermodynamic limit since the continuous spectrum translational Hamiltonian can be obtained as a scaled limit of a sequence of discrete spectrum conformal rotational Hamiltonians associated to global type I systems. In the other hand it has been known for some time that a heat bath chiral KMS state can always be re-interpreted as the Unruh restriction applied to a vacuum system in an larger world (a kind of inverse Unruh effect). Both fact together lead to the above formula for the area density of entropy. In fact using the conformal invariance one can write the area density formula in the more suggestive manner by identifying ε with the conformal invariant cross-ratio of 4 points

$$\varepsilon^2 = \frac{(a_2 - a_1)\,(b_1 - b_2)}{(b_1 - a_1)\,(b_2 - a_2)}$$

where $a_1 < a_2 < b_2 < b_1$ so that (a_1, b_1) corresponds to the larger localization interval and (a_2, b_2) is the approximand which goes with the interpolating type I algebras. At this point one makes contact with some interesting work on what condensed matter physicist call the "entanglement entropy"[10].

One expects that the arguments for the absence of transverse vacuum fluctuations carry over to other null-surfaces as e.g. the upper horizon $\partial\mathcal{D}$ of the double cone $\mathcal{D}$. In the interacting case it is not possible to obtain $\partial\mathcal{D}$ generators through test function restrictions. For zero mass free fields there is however the possibility to conformally transform the wedge into the double cone and in this way obtain the holographic generators as the conformally transformed generators of $\mathcal{A}(\partial W)$. In order to show that the resulting $\mathcal{A}(\partial\mathcal{D})$ continue to play their role even when the bulk generators cease to be conformal one would have to prove that certain double-cone affiliated inclusions are modular inclusions. We hope to return to this interesting problem.

We have presented the pointlike approach and the algebraic approach next to each other, but apart from the free field we have not really connected them. Although one must leave a detailed discussion of their relation to the future, there are some obvious observations one can make. Since for chiral fields the notion of short-distance dimension and rotational spin (the action of the L_0 generator) are closely connected and since the algebraic process of taking relative commutators is bosonic, the lightfront algebras are necessarily bosonic. A field

[10] In [21] the formula for the logarithmically increasing entropy is associated with a field theoretic cutoff and the role of the vacuum polarization cloud in cunjunction with the KMS thermal properties (which is not compatible with a quantum mechanical entanglement interpretation [16]) are not noticed. Since there is no implementation of the split property, the idea of an attenuation of the vacuum polarization cloud has no conceptual place in a path integral formulation. QM and QFT are not distinguished in the functional integral setting and even on a metaphorical level there seems to be no possibility to implement the split property.

as the chiral order variable of the Ising model with dimension $\frac{1}{16}$ does not appear in the algebraic holography but, as mentioned above, it is the pointlike projection of the massive order variable in the factorizing Ising model in the bulk. On the other hand an integer dimensional fields as the stress-energy tensor, is common to both formulations. This suggests that the anomalous dimensional fields which are missing in the algebraic construction may be recovered via representation theory of the transverse extended chiral observable algebra which arises as the image of the algebraic holography.

Since the original purpose of holography similar to that of that of its ill-fated lightcone quantization predecessor, is to achieve a simplified but still rigorous description (for the lightcone quantization the main motivation was a better description of certain "short distance aspects" of QFT), the question arises if one can use holography as a tool in a more ambitious program of classification and construction of QFTs. In this case one must be able to make sense of *inverse holography* i.e. confront the question whether, knowing the local net on the lightfront. one can only obtain at least part of the local substructure of the bulk. It is immediately clear that one construct that part in the bulk which arises from intersecting the LF-affiliated wedge algebras. The full net is only reconstructible if the action of those remaining Poincaré transformations outside the 7-parametric LF covariance group is known.

The presence of the Moebius group acting on the lightlike direction on null-surfaces in curved spacetime resulting from bifurcate Killing horizons [22] has been established in [23], thus paving the way for the transfer of the thermal results to QFT in CST. This is an illustration of symmetry enhancement which is one of holographies "magics".

The above interaction-free case with its chiral abelian current algebra structure (1) admits a much larger unitarily implemented symmetry group, namely the diffeomorphism group of the circle. However the unitary implementers (beyond the Moebius group) do not leave the vacuum invariant (and hence are not Wigner symmetries). As a result of the commutation relations (3) these Diff(S^1) symmetries are expected to appear in the holographic projection of interacting theories. These unitary symmetries act only geometrically on the holographic objects; their action on the bulk (on which they are also well-defined) is fuzzy i.e. not describable in geometric terms. This looks like an interesting extension of the new setting of local covariance [24]

The area proportionality for localization entropy is a structural property of LQP which creates an interesting and hopefully fruitful contrast with Bekenstein's are law [25] for black hole horizons. Bekenstein's thermal reading of the area behavior of a certain quantity in classical Einstein-Hilbert like field theories has been interpreted as being on the interface of QFT with QG. Now we see that the main support, namely the claim that QFT alone cannot explain an area behavior, is not correct. There remains the question whether Bekenstein's numerical value, which people tried to understand in terms of quantum mechanical level occupation, is a credible candidate for quantum entropy. QFT gives a family of area laws with different vacuum polarization *attenuation parameters* ε and it is easy to fix this parameter in terms of the Planck length so that the

two values coalesce. The problem which I have with such an argument is that I have never seen a situation where a *classical* value remained intact after passing to the quantum theory. This does only happen for certain *quasiclassical* values in case the system is integrable.

3 From holography to correspondence: the AdS-CFT correspondence and a controversy

The holography onto null-surfaces addresses the very subtle relation between bulk quantum matter and the projection onto its causal/event horizon as explained in the previous section. A simpler case of holography arises if the bulk and a lower dimensional brane[11] (timelike) boundary share the same maximally possible spacetime (vacuum) symmetry. The only case where this situation arises between two global Lorentz manifolds of different spacetime dimension is the famous AdS-CFT correspondence. In that case the causality leakage off a brane does not occur. In the following we will use the same terminology for the universal coverings of AdS/CFT as for the spacetimes themselves.

Already in the 60s the observation that the 15-parametric conformal symmetry which is shared between the conformal of 3+1-dimensional compactified Minkowski spacetime and the 5-dim. Anti-de-Sitter (the negative constant curvature brother of the cosmologically important de Sitter spacetime) brought a possible field theoretic relation between these theories into the foreground; in fact Fronsdal [26] suspected that QFTs on both spacetimes share more than the spacetime symmetry groups. But the modular localization theory which could convert the shared group symmetry into a relation between two *different spacetime ordering devices* (in the sense of Leibniz) for the *same abstract quantum matter substrate* was not yet in place at that time. Over several decades the main use of the AdS solution has been (similar to Goedel's cosmological model) to show that Einstein-Hilbert field equations besides the many desired solution (as the Robertson-Walker cosmological models and the closely related de Sitter spacetime) also admit unphysical solutions (leading to timelike selfclosing worldlines, time machines, wormholes etc.) and therefore should be further restricted.

The AdS spacetime lost this role of only providing counterexamples and began to play an important role in particle physics when the string theorist placed it into the center of a conjecture about a correspondence between a particular maximally supersymmetric massless conformally covariant Yang-Mills model in d=1+3 and a supersymmetric gravitational model. The first paper was by J. Maldacena [27] who started from a particular compactification of 10-dim. superstring theory, with 5 uncompactified coordinates forming the AdS spacetime. Since the mathematics as well as the conceptual structure of string theory is

[11] In general the brane has a lower dimensional symmetry than its associated bulk and usually denotes d-1 dimensional subspace which contains a time-like direction. Different from null-surfaces branes have a causal leakage.

poorly understood, the string side was identified with one of the supersymmetric gravity models which inspite of its being non-renormalizable admitted a more manageable Lagrangian formulation and was expected to have a similar particle content. On the side of CFT he placed a maximally supersymmetric gauge theory of which calculations which verify the vanishing of the low order beta function already existed[12] (certainly a *necessary* prerequisite for conformal invariance). The arguments involved perturbation theory and additional less controllable approximations. The more than 4.700 follow up papers on this subject did essentially not change the status of the conjecture. But at least some aspects of the general AdS-CFT correspondence became clearer after Witten [28] exemplified the ideas in the field theoretic context of a Φ^4 coupling on AdS using a Euclidean functional integral setting.

The structural properties of the AdS-CFT correspondence came out clearly in Rehren's [30] *algebraic holography.* The setting of local quantum physics (LQP) is particularly suited for questions in which one theory is assumed as given and one wants to construct its holographic projection or its corresponding model on another spacetime. LQP can solve such problems of isomorphisms between models without being forced to actually construct a model on either side (which functional integration proposes to do but only in a metaphoric way) be. At first sight Rehren's setting rewritten in terms of functional integrals (with all the metaphoric caveats, but done in the best tradition of the functional trade) looked quite different from Witten's functional representation. But thanks to a functional identity (explained in the Duetsch-Rehren paper) which shows that fixing functional sources on a boundary and forcing the field values to take on a boundary value via delta function in the functional field space leads to the same result. In this way the apparent disparity disappeared [31] and there is only one AdS-CFT correspondence within QFT.

There are limits to the rigor and validity of functional integral tools in QFT. Even in QM where they are rigorous an attempt to teach a course on QM based on functional integrals would end without having been able to cover the standard material. As an interesting mental exercise just image a scenario with Feynman before Heisenberg. Since path integral representations are much closer to the old quasiclassical Bohr Sommerfeld formulation the transition would have been much smoother, but it would have taken a longer time to get to the operational core of quantum theory; on the other hand quasiclassical fomulas and perturbative corrections thereof would emerge with elegance and efficiency.

Using the measure theoretical functional setting it is well-known that super-renormalizable polynomial couplings can be controlled this way [35]. Realistic models with infinite wave function renormalization constants (all realistic Lagrangian models in more than two spacetime dimensions have a trans canonical short distance behavior) do not fall into this amenable category. But even in low dimension, where there exist models with finite wave function renormalization

[12] An historically interesting case in which the beta function vanishes in every order is the massive Thirring model. In that case the zero mass limit is indeed conformally invariant, but there is no interacting conformal theory for which a perturbation can be formulated directly, it would generate unmanagable infrared divergencies.

constants and hence the short distance prerequisites are met, the functional setting of the AdS-CFT correspondence has an infrared problem[13] of a nasty unresolved kind [37]. As the result of lack of an analog to the operator formulation in QM the suggestive power, their close relation to classical geometric concepts and their formal elegance functional integrals have maintained their dominant role in particle physics although renormalized perturbation theory is better taken care of in the setting of "causal perturbation". An operator approach which is not only capable to establish the mathematical existence of models but also permits their explicit construction exists presently only in d=1+1; it is the previously mentioned bootstrap-formfactor or wedge-localization approach for fsctorizing models. Lagrangian factorizing models only constitute a small fraction.

For structural problems as holography, where one starts from a given theory and wants to construct its intrinsically defined holographic image, the use of metaphorical instruments as Euclidean functional integral representations is suggestive but not really convincing in any mathematical sense. As in the case of lightfront holography there are two mathematically controllable ways to AdS-CFT holography; either using (Wightman) fields (*projective holography*) or using operator algebras (*algebraic holography*). The result of all these different methods can be consistently related [31][32].

The main gain in lightfront holography is a significant simplification of certain properties as compared to the bulk. Even if some of the original problems of the bulk come back in the process of holographic inversion they reappear in the more amenable form of several smaller problems rather than one big one.

The motivation for field theorists being interested in the AdS-CFT correspondence is similar, apart from the fact that the simplification obtainable through an *algebraic isomorphism* is more limited (less radical) than that of a projection. Nevertheless it is not unreasonable to explore the possibility whether some hidden property as e.g. a widespread conjectures *partial integrabilty*[14] could become more visible after a spacetime "re-packaging" of the quantum matter substrate from CFT to AdS.

Despite many interesting analogies between chiral theories and higher dimensional QFT [36] little is known about higher-dimensional conformal QFTs. There are Lagrangian candidates as e.g. certain supersymmetric Yang-Mills theories which fulfill (at least in lowest order) some perturbative prerequisite of conformality which consists in a vanishing beta-function. As mentioned before perturbation theory for conformal QFT, as a result of severe infrared problems, cannot be formulated directly. The prime example for such a situation is the massive Thirring model for which there exists an elegant structural argument for $\beta(g) = 0$ and the knowledge about the non-perturbative massless version

[13] Infrared problems of the kind as they appear in interacting conformal theories are strictly speaking not susceptible to perturbation theoretical treatment and they also seem to pose serious (maybe unsoluble) problems in functional integral representations. In those cases where on knows the exact form of the massless limit (Thirring model) this knowledge can be used to disentangle the perturbative infrared divergences.

[14] Global integrability is only possible in d=1+1, but I am not aware of any theorem which rules out the possibility of integrable substructures.

can then be used to find the correct perturbative infrared treatment.

As far as I could see (with appologies in case of having overlooked some important work) none of these two steps has been carried out for SUSY-YM, so even the conformal side of the Maldacena conjecture has remained unsafe territory.

There is one advantage which null-surface holography has over AdS-CFT type brane holography. The cardinality of degrees of freedom adjusts itself to what is *natural* for null-surfaces (as a manifold in its own right); for the lightfront holography this is the operator algebra generated from extended chiral fields (3). On the other hand this "thinning out" in holographic projections is of course the reason whay inverse holography becomes more complicated and cannot be done with the QFT on one null surface only.

In the holography of the AdS-CFT correspondence the bulk degrees of freedom pass to a conformal brane; in contradistinction to the holography on null-surfaces there is *no reduction of degrees of freedom* resulting from projection. Hence the AdS−CFT isomorphism starting from a "normal" (causally complete as formally arising from Lagrangians) 5-dimensional AdS leads to a *conformal field theory with too many degrees of freedom*. Since a "thinning out" by hand does not seem to be possible, the "physically health" of such a conformal QFT is somewhat dodgy, to put it mildly.

In case one starts with a free Klein-Gordon field on AdS one finds that the generating conformal fields of the CFT are special *generalized free fields* i.e. a kind of continuous superpositions of free fields. They were introduced in the late 50s by W. Greenberg and their useful purpose was (similar to AdS in classical gravity) to *test the physical soundness of axioms of QFT* in the sense that if a system of axioms allowed such solutions, it needed to be further restricted [33] (in that case the so-called causal completion or time-slice property excluded generalized free fields). It seems that meanwhile the word "physical" has changes its meaning, it is used for anything which originated from a physicist.

In the opposite direction the degrees of freedom of a "normal" CFT become "diluted" on AdS in the inverse correspondence. There are not sufficient degrees of freedom for arriving at nontrivial compactly localized operators, the cardinality of degrees of freedom is only sufficient to furnish noncompact regions as AdS wedges with nontrivial operators, the compactly localized double cone algebras remain trivial (multiples of the identity). In the setting based on fields this means that the restriction on testfunction spaces is so severe that pointlike field $A_{AdS}(x)$ at interior points $x \in intAdS$ do not exist in the standard sense as operator-valued distributions on Schwartz spaces. They exist on much smaller test function spaces which contain no functions with compact localizations.

Both sides of the correspondence have been treated in a mathematically rigorous fashion for free AdS (Klein-Gordon equation) theories and free (wave equation) CFT [34][32] where the mismatch between degrees of freedom can be explicated and the structural arguments based on the principles of general QFT show that this mismatch between the transferred and the natural cardinality of the degree of freedom is really there. In terms of the better known Lagrangian formalism the statement would be that if one starts from a Lagrange theory at

one side the other side cannot be Lagrangian. Of course both sides remain QFT in the more general sense of fulfilling the required symmetries, have positive energy and being consistent with spacelike commutativity. In the mentioned free field illustration a AdS Klein-Gordon field is evidently Lagrangian whereas the corresponding *conformal generalized free field* has no Lagrangian and cannot even be characterized in terms of a local hyperbolic field equation. According to the best educated guess, 4-dim. maximally supersymmetric Yang-Mills theories (if they exist and are conformal) would be a natural conformal QFTs "as we know it" and therefore cannot come from a natural QFT on AdS. Needless to say again that there are severe technical problems to set up a perturbation theory for a conformally invariant interactions, the known perturbative systematics breaks down in the presence of infrared problems[15].

I belong to a generation for which not everything which is mathematically possible must have a physical realization; in particular I do not adhere to the new credo that every mathematically consistent idea is realized in some parallel world (anthropic principle): no parallel universe for the physical realization of every mathematical belch.

Generalized free fields[16] and their interacting counterparts which arise from natural AdS free- or interacting- fields remain in my view unphysical, but are of considerable mathematical interest. They do not fit into the standard causal localization setting and they do not allow thermal KMS states without a limiting Hagedorn temperature (both facts are related). Nature did not indicate that it likes to go beyond the usual localizability and thermal behavior. If string theory demands such things it is not my concern, let Max Tegmark find another universe where nature complies with string theory.

Holography is a technical tool and not a physical principle. It simplifies certain aspects of a QFT at the expense of others (i.e. it cannot achieve miracles). The use of such ideas in intermediate steps may have some technical merits, but I do not see any scientific reason to change my viewpoint about physical admissibility. The question of whether by changing the spacetime encoding one could simplify certain properties (e.g. detect integrable substructures) of complicated theories is of course very interesting, but in order to pursue such a line it is not necessary to physically identify the changed theory. Such attempts where only one side needs to be physical and the role of holography would consist in exposing certain structural features which remained hidden in the original formulation sound highly interesting to me.

There is however one deeply worrisome aspect of this whole development. Never before has there been more than 4.700 publication on such a rather narrow subject; in fact even nowadays, one decade after this gold-digger's rush

[15] A well-known problem is the massive Thirring model which leads to $\beta = 0$ in all orders. In this case one already knew confmal limit in closed form and was able to check the correctness of the relation by consistency considerations.

[16] It is interesting to note that the Nambu-Goto Lagrangian (which describes a classical relativistic string) yields upon quantization a pointlike localized generalized free field with the well-known infinite tower mass spectrum and the appearance of a Hagedorn limit temperature. As such it is pointlike localized and there is *no intrinsic quantum concept* which permits to associate it with any stringlike localization.

about the AdS-CFT correspondence started, there is still a sizable number of papers every month by people looking for nuggets at the same place but without bringing Maldacena's gravity-gauge theory conjecture any closer to a resolution. Even with making all the allowances in comparison with earlier fashions, this phenomenon is too overwhelming order to be overlooked. Independent of its significance for particle physics and the way it will end, the understanding of what went on and its covering by the media will be challenging to historians and philosophers of science in the years to come.

I know that it is contra bonos mores to touch on a sociological aspect in a physics paper, but my age permits me to say that at no time before was the scientific production in particle theory that strongly coupled to the Zeitgeist as during the last two decades; never before had global market forces such a decisive impact on the scientific production. Therefore it is natural to look for an explanation why thousands of articles are written on an interesting (but not clearly formulated) conjecture with hundreds of other interesting problems left aside; where does the magic attraction come from? Is it the Holy Grail of a TOE which sets into motion these big caravans? Did the critical power of past particle physics disappear in favor of acclamation? Why are the few critical but unbiased attempts only mentioned by the labels given to them and not by their scientific content?

Since commentaries about the crisis in an area of which one is part run the risk of being misunderstood, let me make perfectly clear that particle physics was a speculative subject and I uphold that it must remain this way. Therefore I have no problem whatsoever with Maldacena's paper; it is in the best tradition of particle physics which was always a delicate blend of a highly imaginative and innovative contribution from one author with profoundly critical analysis of others. I am worried about the loss of this balance. My criticism is also not directed against the thousands of authors who enter this area in good faith believing that they are working at an epoch-forming paradigmatic problem because their peers gave them this impression. Even if they entered for the more mundane reason of carving out a career, I would not consider this as the cause of the present problem.

The real problem is with those who by their scientific qualifications and status are the intellectual leaders and the role models. If they abdicate their role as critical mediators by becoming the whips of the TOE monoculture of particle physics then checks and balances will be lost. Would there have been almost 5000 publication on a rather narrow theme (compared with other topics) in the presence of a more critical attitude from leading particle physicists? No way. Would particle theory, once the pride of theoretical physics with a methodological impact on many adjacent areas have fallen into disrespect and be the object of mock within the larger physics community? The list of questions of this kind with negative answers can be continued.

It is worthwile to look back at times when the delicate balance between the innovative and speculative on the one hand and the critical on the other was still there. Young researchers found guidance by associating themselves to "schools of thought" which where associated with geographical places and names

as Schwinger, Landau, Bogoiubov, Wheeler, Wightman, Lehmann, Haag... who represented different coexisting schools of thought. Instead of scientific cross fertilization between different schools, the new globalized caravan supports the formation of a gigantic monoculture and the loss of the culture of checks and balances.

Not even string theorists can deny that this unfortunate development started with string theory. Every problem string theory addresses takes on a strange metaphoric aspect, an effect which is obviously wanted as the fondness for the use of the letter M shows. The above mentioned AdS-CFT topic gives an illustration which (with a modest amount of mathematical physics) shows the clear structural QFT theorem as compared to the strange conjecture which even thousands of publications were not able to liberate from the metaphoric twilight.

But it is a remarkable fact that, whenever string theorist explain their ideas by QFT analogs in the setting of functional integrals as was done by Witten in [28] for the φ^4 coupling, and on the other hand algebraic quantum field theorists present their rigorous structural method for the same model in the same setting [31], the two results agree (see also [37]).

This is good news. But now comes the bad news. Despite the agreement the Witten camp, i.e. everybody except a few individuals, claim that there exist two different types of AdS-CFT correspondences namely theirs and another one which at least some of them refer to as the "German AdS-CFT correspondence". Why is that? I think I know but I will not write it.

At this point it becomes clear that it is the abandonment of the critical role of the leaders which is fuelling this unhealthy development. Could a statement: "X-Y-Z theory is a gift of the 21st century which by chance fell into the 20 century" have come from Pauli, Schwinger, or Feynman? One would imagine that in those days people had a better awareness that mystifications like this could disturb the delicate critical counterbalance which the speculative nature of particle physics requires. The long range negative effect on particle theory of such a statement is proportional to the prominence and charisma of its author.

There have been several books which criticise string theory. Most critics emphasize that the theory has not predicted a single observable effect and that there is no reason to expect that this will change in the future. Although I sympathize with that criticism, especially if it comes from experimentalists and philosphers, I think that a theorist should focus his critique on the conceptual and mathematical structure and not rely on help from Karl Popper or dwell on the non-existent observational support. Surprisingly I could not find any scholarly article in this direction. One of the reasons may be that after 4 decades of development of string theory such a task requires rather detailed knowledge about its conceptual and mathematical basis. As a result of this unsatidfactory situation I stopped my critical article [29] from going into print and decided to re-write it in such a way that the particle physics part is strengthened at the expense of the sociological sections.

The aforementioned situation of ignoring results which shed a critical light on string theory or the string theorists version of the AdS-CFT correspondence is

perhaps best understood in terms of the proverbial *executing of the messenger who brings bad news*; the unwanted message in the case at hand being the *structural* impossibility to have Lagrangian QFTs with causal propagation on both sides of the correspondence.

It seems that under the corrosive influence of more than 4 decades of string theory, Feynman's observation about its mode of arguing being based on finding excuses instead of explanations, which two decades ago was meant to be provocative, has become the norm. The quantum gravity-gauge theory conjecture is a good example of how a correct but undesired AdS-CFT correspondence is shifted to the elusive level of string theory and quantum gravity so that the degrees of freedom aspect becomes pushed underneath the rug of the elusive string theory where it only insignificantly enlarges the already very high number of metaphors.

There have been an increasing number of papers with titles as "QCD and a Holographic Model of Hadrons", "Early Time Dynamics in Heavy Ion Collisions and AdS/CFT Correspondence", "Confinement/Deconfinement Transition in AdS / CFT", "Isospin Diffusion in Thermal AdS/CFT with flavour", "Holographic Mesons in a Thermal Bath", "Viscous Hyrodynamics and AdS/CFT", "Heavy Quark Diffusion from AdS/CFT".... Ads/CFT for everything? Is string theory bolstered by AdS-CFT really on the way to become a TOE for all of physics, a theory for anything which sacrifies conceptual cohesion to amok running calculations? Or are we witnessing a desperate attempt to overcome the more than 4 decade lasting physical disutiliy? Perhaps it is only a consequence of the "liberating" effect of following prominent leaders who have forgone their duty as critical mediators and preserver of conceptual cohesion.

4 Concluding remarks

In these notes we revisited one of the oldest and still unsolved conceptual problems in QFT, the existence of interacting models. Besides some new concrete results about the existence of factorizing models (which only exist in d=1+1), it is the new method itself, with its promise to explore new fundamental and fully intrinsic properties of QFT, which merits attention. A particularly promising approach for the classification and construction of QFTs consists in using holographic lightfront projections (and in a later stage work one's way back into the bulk). In this situation the holographic degrees of freedom are thinned out as compared to the bulk i.e. the extended chiral fields have lesser number of degrees of freedom.

The concept of degrees of freedom used here is a dynamical one. Knowing only a global algebra[17] as the wedge algebra i.e. $\mathcal{A}(W) \subset B(H)$ as an inclusion into the full algebra one uses fewer degrees freedom than one needs in order to describe the full local substructure of $\mathcal{A}(W)$ i.e. knowing $\mathcal{A}(W)$ in the sense

[17]Knowing an operator algebra means knowing its position within the algebra $B(H)$ of all operators. Knowing its net substructure means knowing the relative position of all its subalgebras.

of a local net. The degrees of freedom emerge always from relations between algebras whereas the single algebra is a structureless monad[15]. Saying that the net $\mathcal{A}(LF)$ has less degrees of freedom than the net associated with the bulk is the same as saying that the knowledge of the LFaffiliated wedges does not suffice to reconstruct the local bulk structure. In this sense the notion of degrees of freedom depends on the knowlege one has about a system; refining the net structure of localized subalgebras of a global algebra increases the degrees of freedom.

The lightfront holography is a genuine projection with a lesser cardinality of degrees of freedom i.e. without knowing how other Poincaré transformations outside the 7-parametric invariance group of the lightfront act it is not uniquely invertible. On its own, i.e. without added information, the lightfront holography cannot distinguish between massive and massless theories; a transverse extended chiral theories does not know whether the bulk was massive or massless. The knowledge of how the opposite lightray translation $U(a_-)$ acts on $\mathcal{A}(LF)$ restores uniqueness; but this action is necessarily "fuzzy" i.e. non-geometric, purely algebraic.Only upon returning to the spacetime ordering device in terms of the bulk it becomes geometric.

The hallmark of null-surface holography is an area law for localization entropy in which the proportionality constant is a product of a holographic matter dependent constant times a logarithmic dependence on the attenuation length for vacuum polarization.

By far the more popular holography has been the AdS-CFT correspondence. Here its physical utility is less clear than the mathematical structure.

Is there really a relation between a special class of conformal gauge invariant gauge theories with supersymmetric quantum gravity? Not a very probable consequence of a change of an spacetime ordering device for a given matter substrate which is what holography means. Integrable substructures within such conformal gauge theories which become more overt on the AdSside? This appears a bit more realistic, but present indications are still very flimsy.

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